

# Nonperturbative scalar-meson resonances with open charm and beauty

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## Abstract

We predict several new bound states and resonances for  $B\pi$ ,  $BK$ , and  $BD$   $S$ -wave elastic scattering, within the same model which was recently employed to explain the  $D_{sJ}^*(2317)^+$  meson. In the charm-nonstrange sector we find a (virtual ?) bound state close to the  $B\pi$  threshold, and moreover two nearby resonances of comparable width in the range 5.9 to 6.1 GeV. We obtain bound states  $B_{s0}^*(5570)$  in  $b\bar{s}$  and  $B_{c0}^*(6490)$  in  $\bar{b}c$ , as well as resonances at higher energies. No  $J^P = 0^+$  states are found at 5.6 GeV (in  $\bar{b}u/d$ ), at 5.7 GeV (in  $b\bar{s}$ ), or at 6.1 GeV (in  $\bar{b}c$ ), where they are predicted from pure confinement.

## 1 Introduction

On studying the vast amount of literature [1, 2] dealing with the phenomena observed in  $S$ -wave meson-meson scattering, one may wonder whether the question “*What is a resonance?*” has any answer at all. Different theoretical considerations even come to opposite conclusions [3, 4, 5, 6, 7] for the same phenomenon.

One reason for the difficulty to unambiguously explain meson-meson-scattering phenomena, may stem from the following. In atomic physics, the perturbative approach towards electromagnetic (EM) transitions is justified, since the EM binding forces are apparently more decisive for the energy-level splittings than photon radiation. We may thus predict the radiation frequencies for atomic transitions to a good accuracy, by just determining the atomic levels of the binding forces. Atomic states have lifetimes which are large as compared with the internal EM frequencies. Hence, electrons complete many revolutions about the nucleus before the system decays. Such systems can be considered stable, like planetary ones. We may, therefore, determine the masses of excited atoms, without making a drastic approximation, just from the EM bound-state spectrum, which procedure is only correct for stable systems. The wave functions for very unstable systems have to be determined by different methods. Most probably, they can be expressed as linear combinations of the bound-state solutions, in which case not just one, but all excitations contribute to the transition amplitude.

In Table 1 we express the stability of quark-antiquark systems by estimating the number of complete cycles of the system before decay. For the internal frequency of  $q\bar{q}$  systems, we take the universal frequency  $\omega$  of the Nijmegen model (see Table 2). We read from Table 1 that pions and Kaons are very stable, whereas the eta mesons are reasonably stable. However, strongly decaying mesons are very unstable.

For example, we are convinced that the structure observed at energies below 1 GeV in  $P$ -wave  $\pi\pi$  scattering, is caused by a quark-antiquark state, the  $\rho$  meson, with a mass of about 770 MeV. But if the  $\rho$  resonance stems from a state described by a quark confined to an antiquark, then that state is far from stable. With a size of approximately 1 fm, its internal frequency must be

Table 1: The ratios of internal frequency and width,  $\omega/\Gamma$ , for some stable and less stable mesons.

meson	width ( $\Gamma$ ) [8]	periods ( $\omega/\Gamma$ )
$\pi^\pm$	$2.53 \times 10^{-8}$ eV	$7.5 \times 10^{15}$
$\pi^0$	7.84 eV	$2.4 \times 10^7$
$K^\pm$	$5.32 \times 10^{-8}$ eV	$3.6 \times 10^{15}$
$K_L$	$1.27 \times 10^{-8}$ eV	$1.5 \times 10^{16}$
$K_S$	$7.37 \times 10^{-6}$ eV	$2.6 \times 10^{13}$
$\eta(547)$	1.2 keV	$1.6 \times 10^5$
$\omega(782)$	8.4 MeV	22.5
$\rho(770)$	149 MeV	1.3
$K^*(892)$	51 MeV	3.7
$\eta'(958)$	0.20 MeV	941

some 200 MeV. However, the lifetime of the  $q\bar{q}$  state is of the same order of magnitude, indicated by its width of 150 MeV. Consequently, the  $q\bar{q}$  system associated with the  $\rho$  resonance barely completes ONE complete revolution during its existence. We may thus wonder whether it makes sense to refer to such a phenomenon as a state, or even a particle, while we just concluded that it does not really exist.

Resonances in scattering are somehow supposed to stem from an underlying spectrum of bound states which are broken up by a weaker interaction. Such a philosophy could in principle work very well for the *stable* mesons  $\pi$  and  $K$ , since their formation is associated with strong interactions, whereas their decays are due to electro-weak processes. Nevertheless, strong decay can have dramatic effects even for the masses of pions and Kaons [9]. Once these effects are fully included in the internal dynamics of pions and Kaons, we may indeed handle electro-weak processes perturbatively [10].

In atomic systems the mechanism for binding is the same as the mechanism for transitions, namely electromagnetism. Nevertheless, since the fine-structure constant  $\alpha$  is small, and transitions are of higher order in  $\alpha$  than the binding forces, one may still approximate EM transitions of atoms by the Born term in a perturbative approach. However, the strong binding force of quarks is of comparable order of magnitude as the strong forces that cause hadronic decay. Consequently, in order to analyse hadronic scattering data, one must treat both aspects of strong interactions at least on the same footing, maybe even attributing more importance to strong decay.

The consequences of such a combined approach are, for the time being, still to be studied in effective models, which, however primitive they are and no matter what details need be addressed

in the future, have already indicated two major properties of hadronic resonances:

1. Upon assuming an underlying spectrum of bare mesonic states with  $q\bar{q}$  degrees of freedom only, resonances in meson-meson scattering come out at energies which are very different from the energy levels of the underlying confinement spectrum. This has been confirmed by the Cornell group, focusing on heavy mesonic resonances of the charmonium and bottomonium systems [11,12], the Nijmegen group, describing, in principle, all mesonic resonances [9,13], and the Lisbon group, in the framework of a chiral quark model [14]. All report mass shifts, due to coupled channels, of up to several hundred MeV, which is comparable in magnitude to the level splittings of hadronic resonances.

2. In experiment, more resonances may be observed than just those stemming from the assumed underlying confinement spectrum. The Nijmegen Model predicts several mesonic resonances [3] and even rather stable states [15] which do have a different origin than the majority of mesonic resonances.

## 2 Resonances

Although resonances have not yet been well defined by us, we turn in this section to the issue of their description. Let us first discuss atomic transitions. When a photon passes through a container filled with a gas of non-excited atoms, then it can be absorbed once its frequency matches some allowed transition from the ground state to an excited state of the atom's spectrum. Ideally, the absorption of photons is measured as a bell-shaped fluctuation of the light beam's intensity, when the light frequency is continuously varied. The reason why a Breit-Wigner (BW) form rather accurately fits the shape of this intensity fluctuation is that the Born term of the interaction dominates the process, and moreover, the excited state is relatively stable. Measured over a larger frequency interval, the absorption curve for photons can be described by a sum of BW-type contributions, each one corresponding to a specific excited state of the atom and its respective life time.

Were the fine-structure constant  $\alpha$  not small, then not only higher-order terms in  $\alpha$  would significantly contribute, but also would the photon noticeably couple to all higher excitations. Hence, the form of the scattering amplitude would then certainly not coincide with BW approximations about each peak of the absorption curve. As a consequence, summing up BW-type amplitudes is not justified for strong interactions. This can easily be demonstrated with a simple model.

In Fig. 1, we display the cross section for harmonic-oscillator confinement coupled to meson-meson scattering. The parameters, i.e., the nonstrange and strange quark masses and the oscillator frequency, are chosen (see Table 2) such that the radial excitations come out at 1.39, 1.77,

Table 2: The relevant Nijmegen-model parameters from Ref. [9] (units are in GeV).

$\omega$	up/down mass ( $m_n$ )	$m_s$	$m_c$	$m_b$
0.19	0.406	0.508	1.562	4.724

2.15, 2.53, ... GeV. In the (elastic) meson-meson scattering channel, we have chosen the meson masses at 0.14 GeV (pion) and 0.5 GeV (Kaon). The intensity of the coupling, parametrised by  $\lambda$ , is varied. For  $\lambda = 0.1$  we are in the small-coupling limit (upper-left picture of Fig. 1), for which a sum of BW amplitudes works reasonably well. We also notice that, for  $\lambda = 0.1$ , all peaks are near the energy levels of the pure harmonic oscillator. But as we increase the coupling, we observe that the cross section has structures which are not of a BW type. For  $\lambda = 0.4$ , we see a dominant structure at 0.8 GeV, which in a more realistic approach [3, 16] can be identified with the  $K_0^*(800)$  resonance. From the results of Fig. 1, we conclude that, for the description of strong

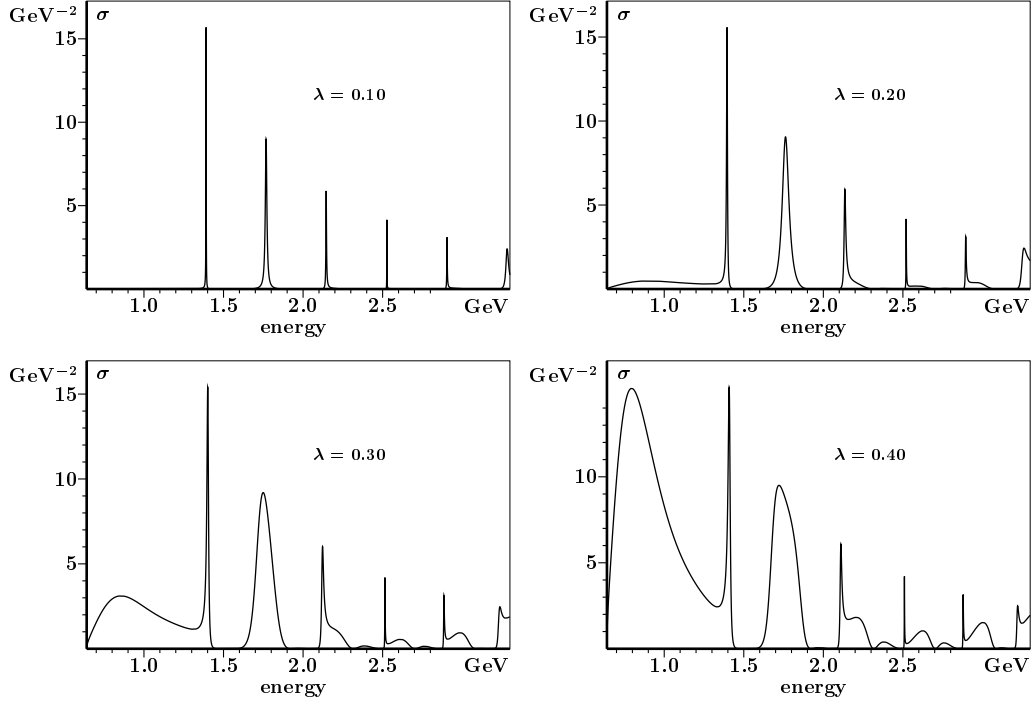


Figure 1: The cross section of elastic  $K\pi$  scattering in the harmonic oscillator model for confinement, for various values of the coupling  $\lambda$ .

interactions, models are needed treating the binding forces and the transition interactions on an equal footing. Our proposal for elastic meson-meson scattering is to use the following expression for the partial-wave phase shift:

$$\cotg(\delta_\ell(p)) = \frac{\pi\lambda^2\mu p \sum_{n=0}^{\infty} \frac{\mathcal{J}_{n\ell}^*(p) \mathcal{N}_{n\ell}(p)}{E(p) - E_{n\ell_c}} - 1}{\pi\lambda^2\mu p \sum_{n=0}^{\infty} \frac{\mathcal{J}_{n\ell}^*(p) \mathcal{J}_{n\ell}(p)}{E(p) - E_{n\ell_c}}} . \quad (1)$$

The details of this formula and its derivation can be found in Refs. [17, 18]. It expresses the fact that in strong hadronic decay all possible confinement states contribute to the scattering amplitude, through the factors  $\mathcal{J}_{n\ell}(p)$  and  $\mathcal{N}_{n\ell}(p)$ , which are convolution integrals over all possible confinement excitations, the transition potential, and spherical Bessel resp. Neumann functions. In the limit of small coupling  $\lambda$ , Eq. (1) reduces to a BW form, with a dominant contribution from the nearest confinement state.

At this stage we may formulate our definition of resonances. Structures like those depicted in Fig. 1 are related to complex-energy singularities in Eq. (1). Those for which a pole in the complex-energy plane (in the nearest Riemann sheet) can be identified are given by us the status of resonance. For small coupling (see Fig. 1 for  $\lambda = 0.1$ ), Eq. (1) has poles near the confinement spectrum, close to the real axis and with small negative imaginary parts. However, poles can

be found, too, with very large (up to several GeVs) imaginary parts, though possibly having no measurable effect. The latter poles can be attributed the status of *mathematical* resonance, but they may be unobservable. Hence, whether or not we include them in our resonance list depends on the accuracy of measurement, not on their actual existence. For stronger coupling we see them showing up as genuine structures in the  $S$ -wave cross section. Suddenly our spectrum has become richer than the original confinement spectrum. This is the realm of strong interactions!

Expression (1), which is the exact solution of a simple dynamical model, is very compact and elegant. But the summations over the radial excitations of the confined system do not converge fast. Hence, for the numerical calculation of scattering quantities it is not suitable. A more practical method has been described in Ref. [19]. Nevertheless, because of its close contact with the physical situation, Eq. (1) is extremely transparent, and useful as a starting point for making approximations.

For example, when we are studying a limited range of energies, we may restrict the summation to only the few terms near by in energy, characterised by  $E_{n\ell_c}$  in Eq. (1), and approximate the remaining terms by a simple function of  $p$ , the relative meson-meson linear momentum. Notice, however, that the resulting expression is only of the type of a BW expansion if the strength of the transition, characterised by  $\lambda$ , is small. For larger values of  $\lambda$ , formula (1) behaves very differently from BW expansions.

Another in practice very useful approximation is inspired by the radial dependence of the transition potential, which was derived in Ref. [20] for meson-meson scattering in the cases  $J^P = 0^-$ ,  $J^P = 1^-$ , and  $J^P = 0^+$ . We find there that it is peaked at short distances. Moreover, if we also take into account that confinement wave functions must be of short range, we may just define a transition radius  $a$  and approximate the spatial integrals in Eq. (1) by choosing a spherical delta shell  $\delta(r - a)$  for the transition potential. Then we end up with the expression

$$\cotg(\delta_\ell(p)) \approx \frac{2a^4\lambda^2\mu p j_\ell(pa) n_\ell(pa) \sum_{n=0}^{\infty} \frac{|\mathcal{F}_{n\ell_c}(a)|^2}{E(p) - E_{n\ell_c}} - 1}{2a^4\lambda^2\mu p j_\ell^2(pa) \sum_{n=0}^{\infty} \frac{|\mathcal{F}_{n\ell_c}(a)|^2}{E(p) - E_{n\ell_c}}} . \quad (2)$$

This way we have pulled the  $p$  dependence outside the infinite summation over the radial confinement spectrum, which makes it much easier to handle truncations, as now the rest term does not depend on  $p$  and may thus be chosen constant.

### 3 Open charm

Our recent detailed analyses of scalar mesons have actually been performed in the approximation of formula (2). For  $DK$  elastic  $S$ -wave scattering, we adopt the same expression that has been used for  $K\pi$  as described in Ref. [16]. We only have to substitute the confinement masses of  $d\bar{s}$  by those for  $c\bar{s}$ , thereby using the parameters of Table 2. For the ground state of the  $J^P = 0^+ c\bar{s}$  confinement spectrum, the mass comes out at 2.545 GeV. The first radial excitation is then 380 MeV (the model's level spacing) higher in mass. The summation over the higher radial excitations in Eq. (2) we approximate by a constant, normalised to 1. The relative couplings of the ground state and the first radial excitation to the  $DK$  channel are 1.0 and 0.2, respectively [16]. Thus we obtain

$$\sum_{n=0}^{\infty} \frac{|\mathcal{F}_{n,1}(a)|^2}{E(p) - E_{n,1}} \longrightarrow \frac{1.0}{E(p) - E_0} + \frac{0.2}{E(p) - E_1} - 1 \quad \text{GeV}^2, \quad (3)$$

where  $E_0 = 2.545$  and  $E_1 = 2.925$ . For Eq. (2) with the substitution (3), and furthermore for the physical value of  $\lambda$  (meaning the value that fits  $K\pi$  elastic  $S$ -wave scattering from threshold up to 1.6 GeV [16]), we show the resulting elastic  $S$ -wave  $DK$  cross section in Ref. [17]. The corresponding pole positions in the complex-energy plane for the  $D_{s0}^*$  ground state and first radial excitation are shown in Ref. [15]. These are found at 2.28 GeV and  $(2.78 - i0.093)$  GeV, respectively. We find no structure in the cross section in the region around 2.55 GeV, where typical quark models predict  $J^P = 0^+ c\bar{s}$  states [21, 12], nor poles in the scattering amplitude. On the contrary, employing confinement *plus* unitarisation, we find structures in the cross section some 250 MeV below and 300–400 MeV above that energy, which explains the observation of the (likely)  $D_{s0}^*$  ground state below the  $DK$  threshold, and predicts for the first radial excitation a resonance somewhere between 2.8 and 3.0 GeV, with a width of roughly 100–200 MeV. We may also predict the scattering length for  $DK$  elastic  $S$ -wave scattering. Keeping in mind that the  $J^P = 0^+ c\bar{s}$  ground state comes out a little bit too low in mass with our value of  $\lambda$ , we estimate the theoretical error by varying  $\lambda$ . For a conservative choice, we obtain for the scattering length in  $DK$   $S$  wave  $a_{DK}^0 = 5 \pm 1 \text{ GeV}^{-1}$ .



## 4 Open beauty

For  $J^P = 0^+$  systems containing one  $b$  quark, we may repeat the above-described procedure. We determine the ground-state masses of the uncoupled systems from the model parameters of Table 2. This results in 5.605, 5.707, and 6.761 GeV for  $\bar{b}u/d$ ,  $b\bar{s}$ , and  $\bar{b}c$ , respectively. The first radial excitations lie 0.38 GeV higher. The remaining parameters  $a$  and  $\lambda$  are kept the same as before, but now scaled with the reduced constituent  $q\bar{q}$  mass  $\mu_{qq}$ , in order to guarantee flavour invariance of the strong interactions [22], i.e.,

$$a_{xy} \sqrt{\mu_{xy}} = \text{constant} \quad \text{and} \quad \lambda_{xy} \sqrt{\mu_{xy}} = \text{constant} \quad , \quad (4)$$

where  $x$  and  $y$  represent the two flavours involved [23]. The constants in Eq. (4) are fixed by [24]  $a_{us} = 3.2 \text{ GeV}^{-1}$  and  $\lambda_{us} = 0.75 \text{ GeV}^{-3/2}$  for  $S$ -wave  $K\pi$  scattering. Henceforth, we shall quote for the parameters  $\lambda$  and  $a$  their values scaled down to elastic  $K\pi$  scattering. The actual values of  $\lambda$  and  $a$  can, in each case, be obtained through Eq. (4) and Table 2.

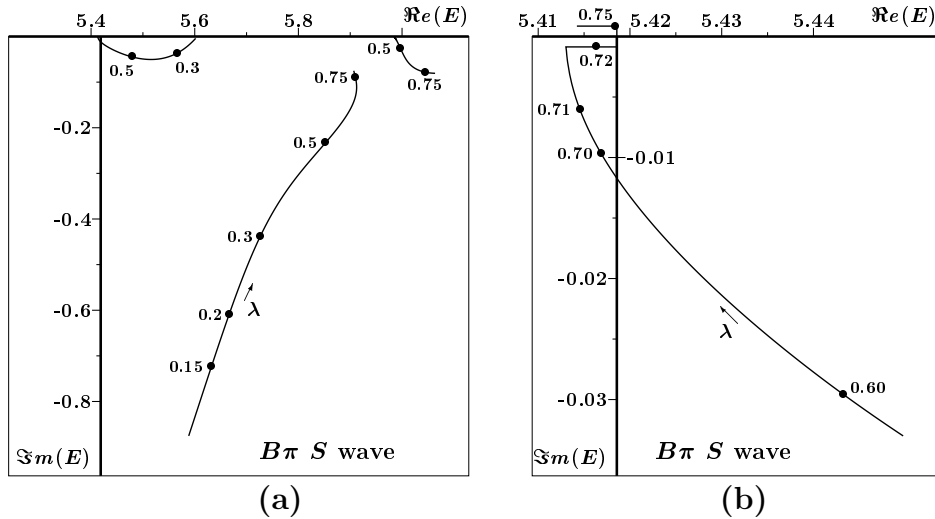


Figure 2: a: Hypothetical pole movements for  $J^P = 0^+$   $B\pi$  elastic scattering. We take  $E_0 = 5605 \text{ MeV}$  and  $E_1 = 5985 \text{ MeV}$ ; b: Details of the lower trajectory.

In Fig. 2 we show how in our model the scattering poles move through the complex-energy plane for  $B\pi$  elastic scattering, when the coupling is varied. For the lowest pole we obtain the very interesting result that the physical pole (for  $\lambda = 0.75$ ) ends up on top of the threshold. Within the accuracy of our model, this means that one has to expect in experiment either a virtual, or a real bound state, just below threshold. Only through an accurate measurement of the scattering length in the  $B\pi$   $S$  wave, the precise position can be obtained experimentally.

For higher energies, we only find poles at energies far above 5.6 GeV. This implies that we do not expect any resonance in the  $S$  wave below roughly 5.9 GeV, where the confinement spectrum

of  $\bar{b}n$  has its first radial excitation. In our model we obtain two resonances, each some 200 MeV wide, and about 150 MeV apart.

In Fig. 3 we depict similar pole trajectories, but now for  $BK$  and  $BD$  elastic scattering. We

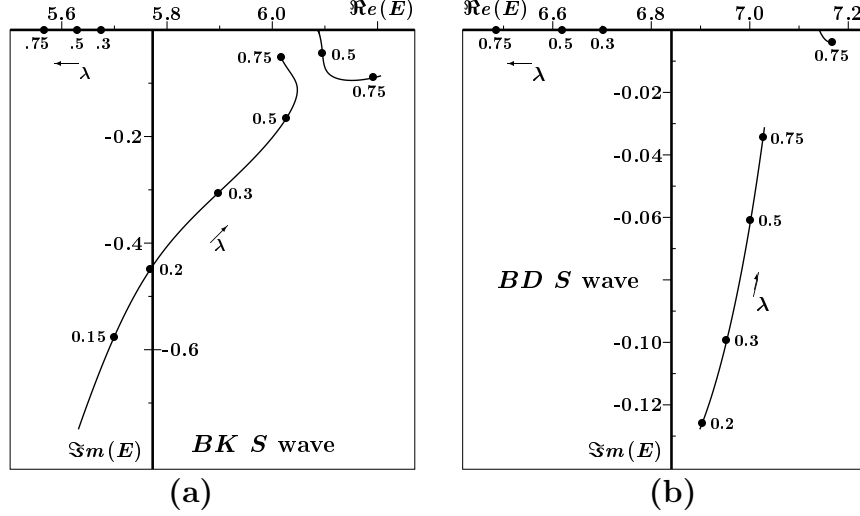


Figure 3: Hypothetical pole movements for  $J^P = 0^+$   $BK$  (a) and  $BD$  (b) elastic scattering. We take  $E_0 = 5707$  MeV and  $E_1 = 6087$  MeV for  $b\bar{s}$  (a), and  $E_0 = 6761$  MeV and  $E_1 = 7141$  MeV for  $\bar{b}c$  (b).

find a  $B_{s0}^*(5570)$  bound state below the  $BK$  threshold, which experimentally should show up as a narrow  $\bar{b}u/d$  resonance. In the  $BD$  case we find a  $B_{c0}^*(6490)$  bound state below threshold, which represents a  $b\bar{s}$  bound state. Also here, nothing is found at 5.707 or 6.761 GeV, where the confinement spectrum has its ground states. Possible resonances above threshold and their widths can be read off from Fig. 3. We summarise these results in Table 3.

Table 3: Bound states and resonances for  $B\pi$ ,  $BK$ , and  $BD$  elastic  $S$ -wave scattering, as read from Figs. 2 and 3.

$q\bar{q}$	ground state mass (GeV)	resonances	
		mass (GeV)	width (GeV)
$\bar{b}u/d$	at $B\pi$ threshold	5.90	0.2
		6.0 – 6.1	0.2
$b\bar{s}$	5.57	6.0 – 6.05	0.1
		6.2	0.2
$\bar{b}c$	6.49	7.0 – 7.05	0.07
		7.18	narrow

## 5 Who is who?

On applying the flavour-independence relations (4) also to  $S$ -wave  $DK$  scattering, we obtained an interesting result, which we shall report here.

In our initial analysis [15], we identified the  $D_{sJ}^*(2317)^+$  meson with the lowest-lying extra pole, which stems from the background, whereas the higher resonance at about 2.9 GeV, with a width of some 150 MeV, was linked to the confinement ground state. However, if we choose for the delta-shell radius  $a$  its corrected flavour-independent value, then we find that the roles of these two poles get interchanged.

This phenomenon is depicted in Fig. 4, where, upon a very small variation in  $a$ , the pole trajectories interchange their physically relevant limiting points, which nevertheless change only very little themselves. As a consequence, we must conclude that the “identity” of a pole cannot

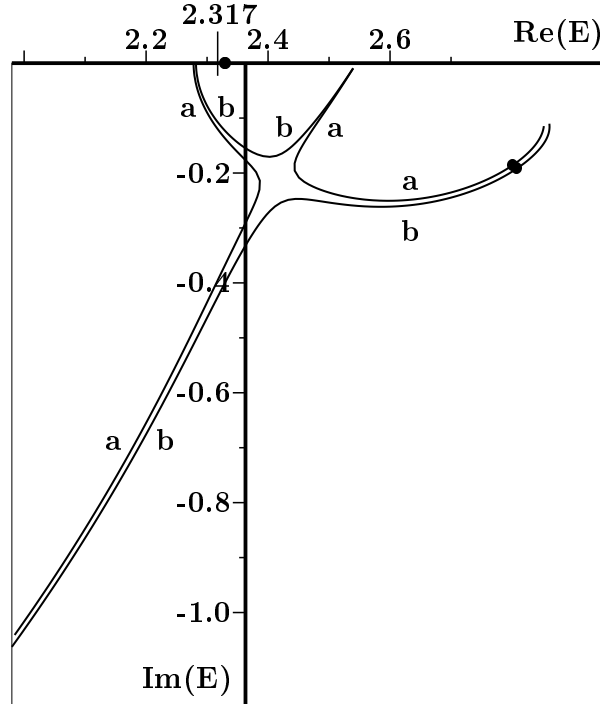


Figure 4: Pole movements for  $DK$   $S$ -wave scattering, with the flavour-invariant radius of the transition potential at the value  $3.45 \text{ GeV}^{-1}$  before crossing (trajectories marked with **a**), and  $3.40 \text{ GeV}^{-1}$  after crossing (trajectories marked with **b**). The  $\bullet$ s correspond to the *new* physical value of  $\lambda$ .

always be inferred with certainty. While in the small-coupling situation (see e.g. the  $\lambda=0.1$  case in Fig. 1) it is unquestionable which poles are to be associated with the confinement spectrum, for intermediate and large couplings one can always find other parameters in Eq. (1) that, when varied simultaneously with the coupling, lead to an identity change. Since intermediate-size couplings

are typical for strong hadronic decay, one must conclude that such processes are inherently nonperturbative, and thus must be treated accordingly.

We should mention, moreover, that the flavour-independence relation (4) mildly alters our previous results. The  $D_{sJ}^*(2317)^+$  pole now comes out at 2.32 GeV (actually very close to the experimental mass), instead of at the published [15] value of 2.28 GeV. Furthermore, we find that the higher resonance shifts even more considerably. Here we obtain a width of some 400 MeV instead of 150 MeV, and a central mass around 2.8 GeV, roughly 100 MeV below the initially proposed value.

## 6 Summary and conclusions

The main issue of this paper is schematically depicted in Fig. 5, where we show how the various sectors of meson physics contribute to resonances in meson-meson scattering. Glue is proba-

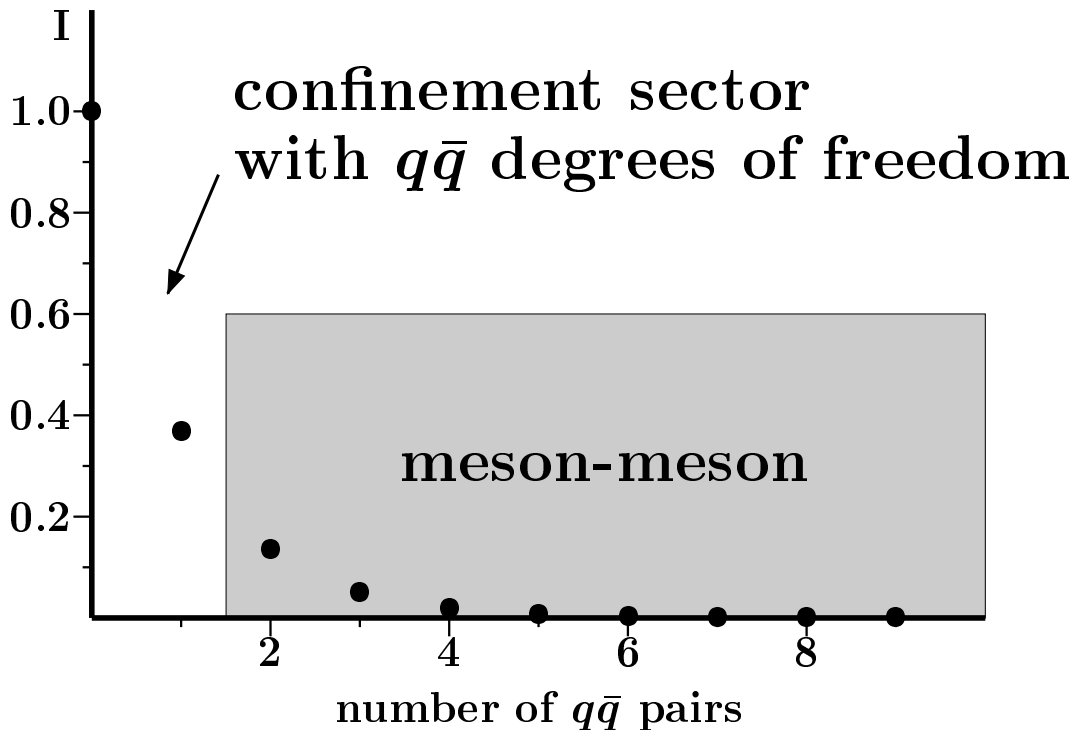


Figure 5: The weight  $I$  of contributions to the dynamics of mesonic states, for different configurations of  $q\bar{q}$  pairs.

bly dominant for the dynamics of meson-meson scattering. It provides confinement, contributes dominantly to the effective masses of the light quarks, and is also involved, though in a complicated and not entirely understood fashion, in the creation of quark-antiquark pairs with vacuum quantum numbers. The degrees of freedom of mesons are given by their  $q\bar{q}$  content. Higher-order configurations of  $q\bar{q}$  pairs mediate the transitions of a mesonic state to multi-meson states. Figure 5 also indicates schematically the origin of deformation of the confinement sector by the multi-meson sector, which results in mass shifts, resonance widths, extra resonances, and bound states. For all that, processes involving the creation of one or more  $q\bar{q}$  pairs are generally strong, and should *not* be handled perturbatively.

As to the spectroscopy of open-charm and open-beauty mesons, we predict several bound states and resonances in  $DK$ ,  $B\pi$ ,  $BK$ , and  $BD$   $S$  waves.

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## References

- [1] J. R. Pelaez, arXiv:hep-ph/0307018; N. G. Kelkar, M. Nowakowski and K. P. Khemchandani, Nucl. Phys. A **724**, 357 (2003) [arXiv:hep-ph/0307184]; T. Kunihiro, arXiv:hep-ph/0307328; S. Ishida and M. Ishida, arXiv:hep-ph/0308307; M. Ishida and S. Ishida, arXiv:hep-ph/0308308; M. Harada, F. Sannino and J. Schechter, arXiv:hep-ph/0309206; M. R. Pennington, presented at *Workshop on Gluonic Excitations, Newport News, Virginia, 14-16 May 2003*, arXiv:hep-ph/0309228; J. R. Pelaez, arXiv:hep-ph/0309292; Pedro Bicudo and Gonalo M. Marques, arXiv:hep-ph/0310008; Muneyuki Ishida and Shin Ishida, arXiv:hep-ph/0310062; Fulvia De Fazio, talk at *International Europhysics Conference on High-Energy Physics (HEP 2003), Aachen, Germany, 17-23 Jul 2003*, arXiv:hep-ph/0310170; M. R. Pennington, arXiv:hep-ph/0310186; K. Yamada, arXiv:hep-ph/0310235; H. Q. Zheng, Z. Y. Zhou, G. Y. Qin, Z. G. Xiao, J. J. Wang and N. Wu, arXiv:hep-ph/0310293; Teiji Kunihiro, Shin Muroya, Atsushi Nakamura, Chiho Nonaka, Motoo Sekiguchi, Hiroaki Wada [SCALAR Collaboration], arXiv:hep-ph/0310312; A. Gillitzer, C. Hanhart, V. Kleber, S. Krewald, H. P. Morsch, F. Rathmann and A. Sibirtsev, arXiv:hep-ph/0311071; A. Gallegos, J. L. Lucio M. and J. Pestieau, arXiv:hep-ph/0311133; W. Ochs, arXiv:hep-ph/0311144.
- [2] E. E. Kolomeitsev and M. F. M. Lutz, arXiv:hep-ph/0311205; J. Hofmann and M. F. M. Lutz, arXiv:hep-ph/0308263; E. E. Kolomeitsev and M. F. M. Lutz, arXiv:hep-ph/0307133; M. F. M. Lutz and E. E. Kolomeitsev, arXiv:nucl-th/0307039; T. E. Browder, S. Pakvasa and A. A. Petrov, arXiv:hep-ph/0307054.
- [3] E. van Beveren, T. A. Rijken, K. Metzger, C. Dullemond, G. Rupp and J. E. Ribeiro, Z. Phys. **C30** (1986) 615.
- [4] T. Hatsuda and T. Kunihiro, Prog. Theor. Phys. Suppl. **91**, 284 (1987).

- [5] Deirdre Black, Amir H. Fariborz, Francesco Sannino, and Joseph Schechter, Phys. Rev. **D58**, 054012 (1998) [hep-ph/9804273].
- [6] S. N. Cherry and M. R. Pennington, Nucl. Phys. A **688** (2001) 823 [arXiv:hep-ph/0005208].
- [7] M. Boglione and M. R. Pennington, Phys. Rev. D **65** (2002) 114010 [arXiv:hep-ph/0203149].
- [8] K. Hagiwara *et al.* [Particle Data Group Collaboration], Phys. Rev. D **66**, 010001 (2002).
- [9] E. van Beveren, G. Rupp, T. A. Rijken, and C. Dullemond, Phys. Rev. D **27**, 1527 (1983).
- [10] A. G. Verschuren, C. Dullemond, and E. van Beveren, Phys. Rev. **D44** (1991) 2803.
- [11] E. Eichten, in *Cargese 1975, Proceedings, Weak and Electromagnetic Interactions At High Energies, Part A*, New York 1976, 305-328.
- [12] M. Di Pierro and E. Eichten, Phys. Rev. D **64**, 114004 (2001) [arXiv:hep-ph/0104208].
- [13] E. van Beveren, C. Dullemond, and G. Rupp, Phys. Rev. **D21** (1980) 772 [Erratum-ibid. **D22** (1980) 787].
- [14] P. J. d. Bicudo and J. E. F. Ribeiro, Phys. Rev. D **42**, 1611 (1990).
- [15] E. van Beveren and G. Rupp, Phys. Rev. Lett. **91**, 012003 (2003) [arXiv: hep-ph/0305035].
- [16] Eef van Beveren and George Rupp, Eur. Phys. J. C **22**, 493 (2001) [arXiv:hep-ex/0106077].
- [17] Eef van Beveren and George Rupp, talk given at *The 25th annual Montreal-Rochester-Syracuse-Toronto Conference on High-Energy Physics*, Joefest, in the honor of the 65th birthday of Joseph Schechter, May 13 - 15, 2003, Syracuse (NY), AIP Conf. Proc. **687**, 86 (2003) [arXiv:hep-ph/0306155].
- [18] E. van Beveren and G. Rupp, Int. J. Theor. Phys. Group Theor. Nonlin. Opt., in press (2003) [arXiv:hep-ph/0304105].
- [19] C. Dullemond, G. Rupp, T. A. Rijken, and E. van Beveren, Comput. Phys. Commun. **27** (1982) 377.
- [20] E. van Beveren, Z. Phys. **C21** (1984) 291.
- [21] Stephen Godfrey and Richard Kokoski, Phys. Rev. D **43**, 1679 (1991).
- [22] K. Abe *et al.* [SLD Collaboration], Phys. Rev. D **59** (1999) 012002 [arXiv:hep-ex/9805023].

- [23] Frieder Kleefeld, contributed to *Hadron 03: 10th International Conference on Hadron Spectroscopy, Aschaffenburg, Germany, 31 Aug - 6 Sep 2003*, arXiv:hep-ph/0310320.
- [24] Eef van Beveren and George Rupp, in *proceedings of 9th International Conference on Hadron Spectroscopy (Hadron 2001), Protvino, Russia, 25 Aug - 1 Sep 2001*, AIP Conf. Proc. **619**, 209 (2002) [arXiv:hep-ph/0110156].